

HOW PENGUINS STARTED TO FLY**The 1999 Sakurai Prize Lecture ^a**

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A mechanism explaining a strong enhancement of nonleptonic weak decays was suggested in 1975, later to be dubbed the penguin. This mechanism extends Wilson's ideas about the operator product expansion at short distances and reveals an intricate interplay of subtle features of the theory such as heavy quark masses in Glashow-Iliopoulos-Maini cancellation, light quarks shaping the chiral properties of QCD, etc. The penguins have subsequently evolved to play a role in a variety of fields in present-day particle phenomenology. I will describe the history of this idea and review its subsequent development. The recent measurement of direct CP violation in K decays gives a new confirmation of the penguin mechanism.

1 History of the idea

It was an exciting period, with Quantum Chromodynamics (QCD) emerging as *the* theory of strong interactions, when three of us – Valya Zakharov, Misha Shifman and I – started in 1973 to work on QCD effects in weak processes. The most dramatic signature of strong interactions in these processes is the so called $\Delta I = 1/2$ rule in nonleptonic weak decays of strange particles. Let me remind you what this rule means by presenting the experimental value for the ratio of the widths of $K_S \rightarrow \pi^+\pi^-$ and $K^+ \rightarrow \pi^+\pi^0$ decays

$$\frac{\Gamma(K_S \rightarrow \pi^+\pi^-)}{\Gamma(K^+ \rightarrow \pi^+\pi^0)} = 450 . \quad (1)$$

The isotopic spin I of hadronic states is changed by $1/2$ in the $K_S \rightarrow \pi^+\pi^-$ weak transition and by $3/2$ in $K^+ \rightarrow \pi^+\pi^0$, so the $\Delta I = 1/2$ dominance is evident.

What does theory predict? The weak interaction has a current \times current form. Based on this, Julian Schwinger suggested¹ to estimate nonleptonic amplitudes as a product of matrix elements of currents, i.e. as a product of semileptonic amplitudes. This approximation, which implies that the strong interaction does not affect the form of the weak nonleptonic interaction, gives

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9/4 for the ratio (1). Thus, the theory is off by a factor of two hundred! We see that strong interactions crucially affect nonleptonic weak transitions.

The conceptual explanation was suggested² by Kenneth Wilson in the context of the Operator Product Expansion (OPE) which he introduced^b. Assuming scaling for OPE coefficients at short distances, Wilson related the enhancement of the $\Delta I = 1/2$ part of the interaction with its more singular behavior at short distances as compared with the $\Delta I = 3/2$ part. In the pre-QCD era it was difficult to test this idea having no real theory of the strong interaction. With the advent of QCD all this changed. The phenomenon of asymptotic freedom gives full theoretical control of short distances. Note in passing that the American Physical Society also followed this development: the discoverers of asymptotic freedom, Gross, Politzer, and Wilczek, became recipients of the 1986 Sakurai Prize.

In QCD the notion of OPE in application to nonleptonic weak interactions can be quantified in more technical terms as one can calculate the effective Hamiltonian for weak transitions at short distances. The weak interactions are carried by W bosons, so the characteristic distances are $\sim 1/m_W$, with $m_W = 80$ GeV. The QCD analysis at these distances in the effective Hamiltonian was done in 1974 by Mary K. Gaillard with Ben Lee⁴, and by Guido Altarelli with Luciano Miani⁵. Asymptotic freedom at short distances means that the strong interaction effects have a logarithmic dependence on momentum rather than power-like behavior, as was assumed in Wilson's original analysis. The theoretical parameter determining the effect is $\log(m_W/\Lambda_{\text{QCD}})$ where Λ_{QCD} is a hadronic scale. These pioneering works brought both good and bad news. The good news was that, indeed, strong interactions at short distances logarithmically enhance the $\Delta I = 1/2$ transitions and suppress the $\Delta I = 3/2$ ones. The bad news was that quantitatively the effect fell short of an explanation of the ratio (1).

Besides $1/m_W$ and $1/\Lambda_{\text{QCD}}$ there are scales provided by masses of heavy quarks t , b and c . In 1975 the object of our study was distances of order $1/m_c$ – the top and bottom quarks were not yet discovered. Introduction of top and bottom quarks practically does not affect the $\Delta S = 1$ nonleptonic transitions. However, it is different with charm. At first sight, the c quark loops looked to be unimportant for nonleptonic decays of strange particles in view of the famous Glashow-Iliopoulos-Miani cancellation⁶ (GIM) with corresponding up quark loops. In 1975 the belief that this cancellation produced the suppression

^bIn Russia we had something of a preview of the OPE ideas worked out by Sasha Patashinsky and Valery Pokrovsky (see their book³) in applications to phase transitions and by Sasha Polyakov in field theories.

factor

$$\frac{m_c^2 - m_u^2}{m_W^2}$$

was universal, which is the reason why the effect of heavy quarks was overlooked. We found instead that:

- (i) The cancellation is distance dependent. Denoting $r = 1/\mu$, we have

$$\begin{aligned} \frac{m_c^2 - m_u^2}{\mu^2}, & \quad \text{for } m_c \ll \mu \leq m_W; \\ \log \frac{m_c^2}{\mu^2}, & \quad \text{for } \mu \ll m_c. \end{aligned} \quad (2)$$

No suppression below m_c !

- (ii) Moreover, new operators appearing in the effective Hamiltonian at distances larger than $1/m_c$ are qualitatively different – they contain right-handed light quark fields in contrast to the purely left-handed structures at distances much smaller than $1/m_c$ (see the next section for their explicit form). It was surprising that right-handed quarks become strongly involved in weak interactions in the Standard Model with its left-handed weak currents. The right-handed quarks are coupled via gluons which carry no isospin; for this reason new operators contribute to $\Delta I = 1/2$ transitions only.
- (iii) For the mechanism we suggested it was crucial that the matrix elements of novel operators were much larger than those for purely left-handed operators. The enhancement appears via the ratio

$$\frac{m_\pi^2}{m_u + m_d} \sim 2 \text{ GeV},$$

which is large due to the small light quark masses. The small values of these masses was a new idea at the time, advocated in 1974 by Heiri Leutwyler⁷ and Murray Gell-Mann⁸. The origin of this large scale is not clear to this day but it shows that in the world of light hadrons there is, besides the evident momentum scale Λ_{QCD} , some other scale, numerically much larger.

Thus, the explanation of the $\Delta I = 1/2$ enhancement comes as a nontrivial interplay of OPE, GIM cancellation, the heavy quark scale, and different

intrinsic scales in light hadrons. I will discuss the construction in more detail below. However, I will first digress to explain what the mechanism we suggested has in common with penguins.

We had a hard time communicating our idea to the world. Our first publication was a short letter published on July 20, 1975 in the Letters to the Journal of Theoretical and Experimental Physics. Although an English translation of JETP Letters was available in the West we sent a more detailed version to Nuclear Physics shortly after. What happened then was a long fight for publication; we answered numerous referee reports – our paper was considered by quite a number of experts. The main obstacle for referees was to overcome their conviction about the GIM suppression $(m_c^2 - m_u^2)/m_W^2$ and to realize that there is no such suppression at distances larger than $1/m_c$. Probably our presentation was too concise for them to follow.

Eventually the paper was published in the March 1977 issue of Nuclear Physics without any revision, but only after we appealed to David Gross who was then on the editorial board. The process took more than a year and a half! We were so exhausted by this fight that we decided to send our next publication containing a detailed theory of nonleptonic decays to Soviet Physics JETP instead of an international journal. Lev Okun negotiated a special deal for us with the editor Evgenii Lifshitz to submit the paper of a size almost twice the existing limit in the journal – paper was in short supply as was almost everything in the USSR. We paid our price: the paper was published in 1977, but even years later many theorists referred to our preprints of the paper without mentioning the journal publication.

On a personal note, let me mention that a significant part of the JETP paper was done over the phone line – Valya and Misha worked in ITEP, Moscow, while I was in the Budker Institute of Nuclear Physics, Novosibirsk. The phone connection was not very good and we paid terrific phone bills out of our own pockets.

I think, that in recognition of our work in the world at large it was Mary K. Gaillard who first broke the ice – she mentioned the idea in one of her review talks. Moreover, she collaborated with John Ellis, Dimitri Nanopoulos, and Serge Rudaz in the work¹⁰ in which they applied a similar mechanism to B physics. It is in this work that the mechanism was christened the penguin.

How come? Figure 1 shows the key Feynman diagram for the new operators in the form we drew it in our original publications⁹. It does not look at all penguin-like, right? Now look how a similar diagram is drawn in the paper of the four authors mentioned above.

You see that that some measures were taken to make the diagram reminiscent of

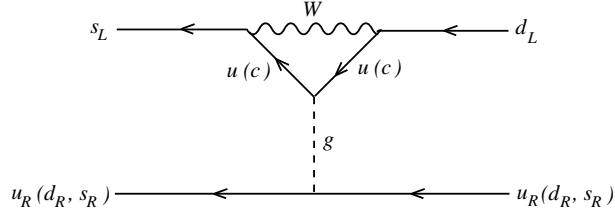


Figure 1: Appearance of new operators due to quark loops

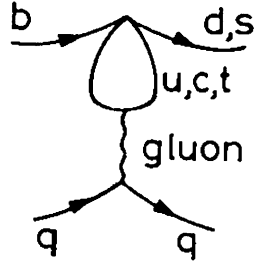


Figure 2: Quark loops in B decays

a penguin. Let me refer here to John Ellis’ recollections^c on how it happened:

“ Mary K, Dimitri and I first got interested in what are now called penguin diagrams while we were studying CP violation in the Standard Model in 1976 ... The penguin name came in 1977, as follows.

In the spring of 1977, Mike Chanowitz, Mary K and I wrote a paper on GUTs predicting the b quark mass before it was found. When it was found a few weeks later, Mary K, Dimitri, Serge Rudaz and I immediately started working on its phenomenology. That summer, there was a student at CERN, Melissa Franklin who is now an experimentalist at Harvard. One evening, she, I and Serge went to a pub, and she and I started a game of darts. We made a bet that if I lost I had to put the word penguin into my next paper. She actually left the darts game before the end, and was replaced by Serge, who beat me. Nevertheless, I felt obligated to carry out the conditions of the bet.

^cJohn sent his recollections to Misha Shifman in 1995, who published them in the preface¹¹ to his book.

For some time, it was not clear to me how to get the word into this b quark paper that we were writing at the time. Then, one evening, after working at CERN, I stopped on my way back to my apartment to visit some friends living in Meyrin, where I smoked some illegal substance. Later, when I got back to my apartment and continued working on our paper, I had a sudden flash that the famous diagrams look like penguins. So we put the name into our paper, and the rest, as they say, is history.”

I learned some extra details of the story from Serge Rudaz who is my Minnesota colleague now. He recollects that for him to beat John in darts was a miraculous event. John was a very strong player and had his own set of darts which he brought to the pub.

2 Effective Hamiltonian

Application of Wilson’s OPE to nonleptonic $\Delta S = 1$ decays means the construction of the effective Hamiltonian H^{eff} as a sum over local operators \mathcal{O}_i ,

$$H^{\text{eff}}(\mu) = \sqrt{2} G_F V_{us}^* V_{ud} \sum_i c_i(\mu) \mathcal{O}_i(\mu). \quad (3)$$

Here V_{us} , V_{ud} are elements of the Cabibbo-Kobayashi-Maskawa mixing matrix and μ denotes the so called normalization point, which is the inverse of the shortest distance for which the effective Hamiltonian is to be applied, \mathcal{O}_i are gauge invariant local operators made out of quark and gluon fields, and c_i are OPE coefficients (c numbers).

At distances larger than $1/m_c$, i.e. at $\mu < m_c$, only light u, d, s quarks and gluons remain as building material for the \mathcal{O}_i . Operators can be ordered according to their canonical dimension d , and the corresponding OPE coefficients are proportional to $(1/m_W)^{d-4}$. As a selection criterion for operators we used their transformation features in the limit of chiral $\text{SU}(2)_L \times \text{SU}(2)_R$ symmetry, picking up operators which are $\text{SU}(2)_R$ singlets. Under this criterion, the operator of lowest dimension ($d = 5$) is of gluomagnetic type (magnetic penguins),

$$T = i \bar{s}_R \sigma_{\mu\nu} t^a d_L G_{\mu\nu}^a, \quad (4)$$

where $G_{\mu\nu}^a$ is the gluon field strength tensor, and t^a ($a = 1, \dots, 8$) are 3×3 generators of the color $\text{SU}(3)$. The corresponding OPE coefficient is proportional to the strange quark mass m_s . For this reason the magnetic penguins turn out to be not important in $\Delta S = 1$ transitions. They are important, however, for the b quark whose mass is large. We will return to this point later.

The operator basis of $SU(3)_R$ invariant operators of $d = 6$ consists of six four-fermion operators. The first four operators are constructed from left-handed quarks (and their antiparticles, which are right-handed),

$$\begin{aligned}
\mathcal{O}_1 &= \bar{s}_L \gamma_\mu d_L \bar{u}_L \gamma^\mu u_L - \bar{s}_L \gamma_\mu u_L \bar{u}_L \gamma^\mu d_L, & (\mathbf{8}_f, \Delta I = 1/2), \\
\mathcal{O}_2 &= \bar{s}_L \gamma_\mu d_L \bar{u}_L \gamma^\mu u_L + \bar{s}_L \gamma_\mu u_L \bar{u}_L \gamma^\mu d_L + 2\bar{s}_L \gamma_\mu d_L \bar{d}_L \gamma^\mu d_L \\
&\quad + 2\bar{s}_L \gamma_\mu d_L \bar{s}_L \gamma^\mu s_L, & (\mathbf{8}_d, \Delta I = 1/2), \\
\mathcal{O}_3 &= \bar{s}_L \gamma_\mu d_L \bar{u}_L \gamma^\mu u_L + \bar{s}_L \gamma_\mu u_L \bar{u}_L \gamma^\mu d_L + 2\bar{s}_L \gamma_\mu d_L \bar{d}_L \gamma^\mu d_L \\
&\quad - 3\bar{s}_L \gamma_\mu d_L \bar{s}_L \gamma^\mu s_L, & (\mathbf{27}, \Delta I = 1/2), \\
\mathcal{O}_4 &= \bar{s}_L \gamma_\mu d_L \bar{u}_L \gamma^\mu u_L + \bar{s}_L \gamma_\mu u_L \bar{u}_L \gamma^\mu d_L \\
&\quad - \bar{s}_L \gamma_\mu d_L \bar{d}_L \gamma^\mu d_L, & (\mathbf{27}, \Delta I = 3/2). \quad (5)
\end{aligned}$$

Every quark field is a color triplet q^i and summation over color indices is implied, $\bar{q}_2 \gamma^\mu q_1 = (\bar{q}_2)_i \gamma^\mu (q_1)^i$. What is marked in the brackets are the $SU(3)$ and isospin features of the operators.

Two more four-fermion operators entering the set contain also right-handed quarks (in $SU(3)_R$ singlet form),

$$\begin{aligned}
\mathcal{O}_5 &= \bar{s}_L \gamma_\mu t^a d_L (\bar{u}_R \gamma^\mu t^a u_R + \bar{d}_R \gamma^\mu t^a d_R + \bar{s}_R \gamma^\mu t^a s_R), \quad (\mathbf{8}, \Delta I = 1/2), \\
\mathcal{O}_6 &= \bar{s}_L \gamma_\mu d_L (\bar{u}_R \gamma^\mu u_R + \bar{d}_R \gamma^\mu d_R + \bar{s}_R \gamma^\mu s_R), \quad (\mathbf{8}, \Delta I = 1/2). \quad (6)
\end{aligned}$$

Operators \mathcal{O}_5 and \mathcal{O}_6 are different by color flow only.

The operator basis we have introduced is recognized now (some doubts were expressed in the literature at the beginning). The standard set¹² used presents some linear combinations of \mathcal{O}_{1-6} . Actually, the set is five instead of six combinations, the completeness of the basis was lost on the way, although it is not important within the Standard Model.

2.1 Evolution

The effective Hamiltonian (2) may remind the reader of the Fermi theory of beta decay with its numerous variants for four-fermion operators. While in many respects the analogy makes sense, the difference is that the standard model together with QCD allows us to fix all coefficients c_i . In the leading logarithmic approximation the evolution of the effective Hamiltonian at $\mu > m_c$ was found in Refs.^{4,5}. Penguins do not appear in this range and the result for

$H^{\text{eff}}(m_c)$ has a simple form:

$$\begin{pmatrix} c_1(m_c) \\ c_2(m_c) \\ c_3(m_c) \\ c_4(m_c) \\ c_5(m_c) \\ c_6(m_c) \end{pmatrix} = \begin{pmatrix} -\left[\frac{\alpha_S(m_b)}{\alpha_S(m_W)}\right]^{4/b_5} \left[\frac{\alpha_S(m_c)}{\alpha_S(m_b)}\right]^{4/b_4} \\ \frac{1}{5} \left[\frac{\alpha_S(m_b)}{\alpha_S(m_W)}\right]^{-2/b_5} \left[\frac{\alpha_S(m_c)}{\alpha_S(m_b)}\right]^{-2/b_4} \\ \frac{2}{15} \left[\frac{\alpha_S(m_b)}{\alpha_S(m_W)}\right]^{-2/b_5} \left[\frac{\alpha_S(m_c)}{\alpha_S(m_b)}\right]^{-2/b_4} \\ \frac{2}{3} \left[\frac{\alpha_S(m_b)}{\alpha_S(m_W)}\right]^{-2/b_5} \left[\frac{\alpha_S(m_c)}{\alpha_S(m_b)}\right]^{-2/b_4} \\ 0 \\ 0 \end{pmatrix}, \quad (7)$$

where $\alpha_S(\mu)$ is the running coupling

$$\alpha_S(\mu) = \frac{\alpha_S(\mu_0)}{1 + b_N \frac{\alpha_S(\mu_0)}{2\pi} \ln \frac{\mu}{\mu_0}}, \quad b_N = 11 - \frac{2}{3}N \quad (8)$$

in the range with N “active” flavors. The modification due to the b quark is not significant, a few percent numerically, and the t quark effects do not appear in the leading logarithmic approximation.

Penguin operators $\mathcal{O}_{5,6}$ show up due to evolution at μ below m_c ,

$$c_i(\mu) = \left[\exp \left\{ \frac{\rho}{b_3} \ln \frac{\alpha_S(\mu)}{\alpha_S(m_c)} \right\} \right]_{ij} c_j(m_c), \quad (9)$$

where the anomalous dimension matrix ρ is

$$\rho = \begin{pmatrix} 34/9 & 10/9 & 0 & 0 & 4/3 & 0 \\ 1/9 & -23/9 & 0 & 0 & -2/3 & 0 \\ 0 & 0 & -2 & 0 & 0 & 0 \\ 0 & 0 & 0 & -2 & 0 & 0 \\ 1/6 & -5/6 & 0 & 0 & 6 & 3/2 \\ 0 & 0 & 0 & 0 & 16/3 & 0 \end{pmatrix}. \quad (10)$$

Numerical results for the OPE coefficients depend on the normalization point μ , pushing μ as low as possible maximizes the effect of the evolution. We chose the lowest μ as the point where $\alpha_S(\mu) = 1$. The values of $c_{1,2,3,4}$ are relatively stable, say, under variation of the c quark mass but the penguin coefficients $c_{5,6}$ depend on it rather strongly. This is not surprising, of course, since penguins are generated in the interval of virtual momenta between μ and

m_c . Numerically, even for μ as low as 200 MeV, the coefficients $c_{5,6}$ are rather small. Our 1975 estimates for c_5 were in the interval

$$c_5 = 0.06 \div 0.14. \quad (11)$$

With the present day coupling $\alpha_S(m_Z) = 0.115$, it would be about twice smaller. The values of the coefficients $c_{5,6}$ are small and unstable. We will return to the problem of OPE coefficients at the low normalization point in connection with the procedure of calculating matrix elements, essentially it is about matching for perturbative and nonperturbative effects.

We also found the coefficient for the magnetic penguin operator (4). Two-loop calculations were necessary for the purpose. Nowadays due to the efforts of few groups the next to leading approximation has been found for all types of weak processes, see, e.g., the review¹³. Although it is nice to have an accurate values of the OPE coefficients the main phenomenological effects come from matrix elements as we will see in the next Section. The theoretical uncertainty is much bigger there.

3 Matrix elements and phenomenology of nonleptonic decays

We used the naive quark model to find matrix elements of four-fermion operators \mathcal{O}_{1-6} . The model implies a factorization for amplitudes of K mesons decays, in the hyperon decays it is also the case for all operators but \mathcal{O}_1 . It is clear that the factorization does not work at the range of μ where the strong coupling $\alpha_S(\mu)$ is small being in contradiction with a calculable evolution. For this reason, if the naive quark model works at all it is only at very low values of μ where the evolution is complete, i.e. where $\alpha_S(\mu) \sim 1$. But then the theoretical accuracy of the perturbative OPE coefficients is not good because it is governed by the same $\alpha_S(\mu)$.

So we employed the following strategy: let us use the naive quark model for hadrons together with a factorization somewhere at small μ but let us allow adjustment of the OPE coefficients from phenomenological fits. It does not involve many parameters. Predominantly three coefficients, c_1 , c_4 and c_5 , plus a few nonfactorizable matrix elements in hyperon decays determine the bulk of nonleptonic amplitudes. New relations arising from such a fit are in good agreement with experimental data.

3.1 $\Delta I = 3/2$ transitions

Let us start with the decay $K^+ \rightarrow \pi^+\pi^0$. It is a $\Delta I = 3/2$ transition and its amplitude is determined by the matrix element of \mathcal{O}_4 ,

$$M(K^+ \rightarrow \pi^+\pi^0) = \sqrt{2} G_F V_{us}^* V_{ud} \langle \pi^+\pi^0 | c_4 \mathcal{O}_4 | K^+ \rangle. \quad (12)$$

In the valence quark model the factorization of this matrix element is visible from the Feynman diagrams presented in Fig. 3.

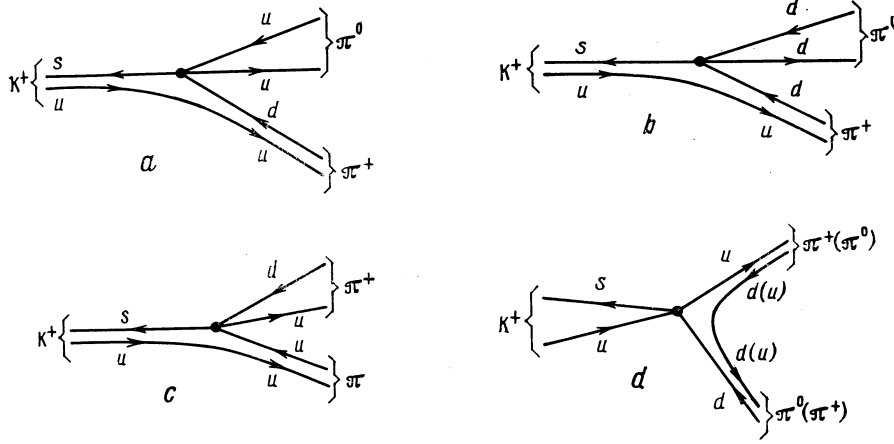


Figure 3: Quark diagrams for K meson decays

Consider, for instance, the diagram a ,

$$\begin{aligned}
 M_4^a &= \langle \pi^0 | \bar{u}_L \gamma_\mu u_L | 0 \rangle \langle \pi^+ | \bar{s}_L \gamma^\mu d_L | K^+ \rangle + \langle \pi^0 | (\bar{u}_L)_i \gamma_\mu (u_L)^j | 0 \rangle \langle \pi^+ | (\bar{s}_L)_j \gamma^\mu (d_L)^i | K^+ \rangle \\
 &= \frac{4}{3} \langle \pi^0 | \bar{u}_L \gamma_\mu u_L | 0 \rangle \langle \pi^+ | \bar{s}_L \gamma^\mu d_L | K^+ \rangle, \quad (13)
 \end{aligned}$$

where out of three terms entering the definition of \mathcal{O}_4 , the first one factorizes, the second factorizes after Fierz transformation, and the third does not contribute. The matrix elements in Eq. (13) are known from semileptonic $\pi \rightarrow \mu \nu$ and $K \rightarrow \pi e \nu$ transitions,

$$\begin{aligned}
 \langle \pi^0 | \bar{u}_L \gamma_\mu u_L | 0 \rangle &= -\frac{i f_\pi}{2\sqrt{2}} q_\mu, \quad f_\pi = 0.95 m_\pi, \\
 \langle \pi^+ | \bar{s}_L \gamma_\mu d_L | K^+ \rangle &= -\frac{1}{2} [(p + q)_\mu f_+ + (p - q)_\mu f_-]. \quad (14)
 \end{aligned}$$

Accounting for all diagrams in Fig. 3 in a similar way we get

$$M(K^+ \rightarrow \pi^+ \pi^0) = i c_4 G_F V_{us}^* V_{ud} m_K^2 f_\pi. \quad (15)$$

Comparing it with the experimental value

$$|M(K^+ \rightarrow \pi^+ \pi^0)|_{\text{exp}} = 0.05 G_F m_K^2 m_\pi \quad (16)$$

we find

$$c_4 \approx 0.25, \quad (17)$$

what is about 1.6 times less than the theoretical estimate of c_4 .

The consistency check comes from the $\Delta I = 3/2$ hyperon decays. In Fig. 4 quark diagrams for the $\Lambda \rightarrow p\pi^-$ decay are presented. The symmetry of wave

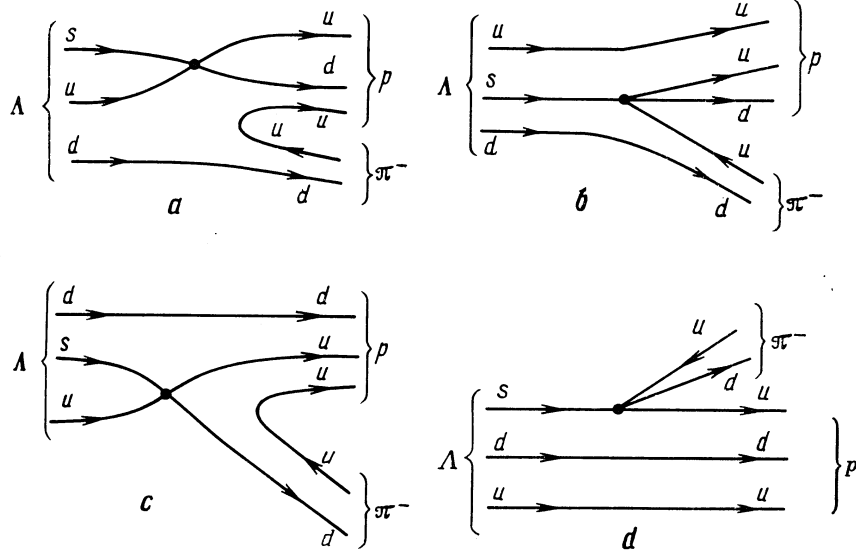


Figure 4: Quark diagrams for hyperon decays

functions and operators under permutations of color indices (first discussed by Pati and Woo¹⁴) is important for the analysis, Namely, the operator \mathcal{O}_1 is antisymmetric under permutation of quark color indices, operators $\mathcal{O}_{2,3,4}$ are symmetric, and operators $\mathcal{O}_{5,6}$ have no specific symmetry, in these quarks differ by their helicities. Baryon wave functions are antisymmetric in color, so only antisymmetric operators survive in diagrams a, b, c . For the symmetric $\Delta I = 3/2$ operator \mathcal{O}_4 only the diagram c remains. For this diagram factorization takes place,

$$\begin{aligned} \langle \pi^- p | \mathcal{O}_4^\dagger | \Lambda \rangle &= \frac{1}{\sqrt{2}} \langle \pi^0 n | \mathcal{O}_4 | \Lambda \rangle = \frac{4}{3} \langle \pi^- | \bar{d}_L \gamma_\mu u_L | 0 \rangle \langle p | \bar{u}_L \gamma^\mu s_L | \Lambda \rangle \\ &\approx -\frac{i}{\sqrt{6}} f_\pi q_\mu \bar{u}_p \left(\gamma^\mu + \frac{5}{9} g_A \gamma^\mu \gamma_5 \right) u_\Lambda. \end{aligned} \quad (18)$$

Using the value (17) for the coefficient c_4 we get predictions for the $\Delta I = 3/2$ hyperon decay amplitudes. They are collected in the Table 1, where the s and p wave amplitudes A and B are defined by

$$M = -iG_F m_\pi^2 \bar{u}_f (A + B\gamma_5) u_i \quad (19)$$

We see that the predictions for s waves are in a reasonable agreement with

Decay	A_{theor}	A_{exper}	B_{theor}	B_{exper}
$\Lambda_-^0 + \sqrt{2}\Lambda_0^0$	0.12	0.09 ± 0.03	-0.95	0.66 ± 0.81
$\sqrt{2}\Sigma_0^+ - \Sigma_+^+ + \Sigma_-^-$	0.14	0.22 ± 0.09	0.49	2.7 ± 1.1
$\Xi_-^- + \sqrt{2}\Xi_0^0$	-0.14	-0.15 ± 0.07	0.23	1.7 ± 2.2

Table 1: $\Delta I = 3/2$ amplitudes in hyperon decays

the data, for p waves the experimental accuracy is too low for a conclusion.

3.2 $\Delta I = 1/2$ transitions

The above analysis of color symmetry and factorization in quark diagrams in Fig. 4 shows that the operators \mathcal{O}_1^\dagger (diagrams a , b , and c) and \mathcal{O}_5^\dagger (the diagram d) are dominant in hyperon decays. Matrix elements of \mathcal{O}_5^\dagger are factorizable but that of \mathcal{O}_1^\dagger are not.

There exists a combination of amplitudes,

$$M(\Xi^- \rightarrow \Sigma^- \pi^0) = \sqrt{3}\Lambda_-^0 - \Sigma_0^+ + \sqrt{3}\Xi_-^-, \quad (20)$$

which does not contain \mathcal{O}_1^\dagger . Thus, the ratio of p and s waves for this combination is predicted,

$$\frac{B(\Xi^- \rightarrow \Sigma^- \pi^0)}{A(\Xi^- \rightarrow \Sigma^- \pi^0)} = g_A \frac{m_\Xi + m_\Sigma}{m_\Xi - m_\Sigma} \approx 25. \quad (21)$$

Experimentally this value is 33 ± 10 . The uncertainty can be reduced if the Lee-Sugawara relation $2\Xi_-^+ + \Lambda_-^+ - \sqrt{3}\Sigma_0^+ = 0$ is used. Then the value is 27.4 ± 2.5 . Thus, we see an experimental confirmation of our description.

The comparison of s wave A

$$A(\Xi^- \rightarrow \Sigma^- \pi^0) = \left(c_5 + \frac{3}{16}c_6\right) V_{us} V_{ud}^* \frac{2}{9} \frac{f_\pi m_\pi^2}{m_u + m_d} \frac{m_\Xi - m_\Sigma}{m_s - m_u} \quad (22)$$

with the corresponding experimental value -0.51 ± 0.10 fixes

$$c_5 + \frac{3}{16}c_6 \approx -0.25, \quad (23)$$

which is about four times larger than the theoretical estimate.

The operators $\mathcal{O}_{5,6}$ are dominant in K decays. The naive quark model, together with σ meson ($m_\sigma \approx 700$ MeV) dominance in the s wave $\pi\pi$ channel, gives for the matrix element

$$\begin{aligned} \langle \pi^+ \pi^- | \mathcal{O}_5 | K_S \rangle &= i \frac{2\sqrt{2}}{9} \frac{f_\pi m_K^2 m_\pi^2}{(m_u + m_d) m_s} \left[\frac{f_K}{f_\pi} \frac{1}{m_\sigma^2 - (q_+ + q_-)^2} - 1 \right] \\ &\approx i \frac{2\sqrt{2}}{9} \frac{f_\pi m_K^2 m_\pi^2}{(m_u + m_d) m_s} \left[\frac{f_K}{f_\pi} - 1 + \frac{f_K}{f_\pi} \frac{m_K^2}{m_\sigma^2} \right]. \end{aligned} \quad (24)$$

The result for the amplitude of $K_S \rightarrow \pi^+ \pi^-$ decay

$$\langle \pi^+ \pi^- | H^{\text{eff}} | K_S \rangle = (0.85 + 0.20) i G_F m_K^2 m_\pi, \quad (25)$$

where the first number comes from $\mathcal{O}_{5,6}$ and the second from \mathcal{O}_{1-4} , matches the experimental value.

4 Further developments

4.1 Limit of large N_{color} , chiral loops

As we see in the approach presented above the main uncertainty comes from the range of momenta of the order of hadronic scales. To make a consistent analysis in this range Bardeen, Buras, and Gerard suggested to use a description in terms of chiral meson dynamics¹⁵. It means that QCD is used to fix the effective Hamiltonian at μ below charm threshold but large enough to have $\alpha_S(\mu) \ll 1$, particularly, they choose $\mu \sim 1$ GeV. Matrix elements of this effective Hamiltonian which accounts for momenta below 1 GeV are calculated with mesonic loops instead of quark loops. The use of chiral meson loops can be parametrically justified in the limit of large number of colors N_c .

Notice that penguin operators appear due to mixing with the left-handed operators \mathcal{O}_{1-4} , this mixing is suppressed as $1/N_c$. In the dual mesonic picture it is hadronic vertices which bring in $1/N_c$. The smallness of $1/N_c$ emphasizes once more the necessity of a large hadronic scale in the problem, it was m_π^2/m_q in the naive quark model. In chiral dynamics $\Delta I = 1/2$ amplitudes also come out enhanced what confirm the penguin explanations of $\Delta I = 1/2$. The chiral loops also improve predictions for $\Delta I = 3/2$ transitions providing an additional suppression.

4.2 Penguin decays of B mesons

We mentioned that “magnetic” penguins, i.e. the $d = 5$ operators of the kind given by Eq. (4) are of particular importance for B decays (see review¹⁶). The Feynman diagram for $b \rightarrow sg(\gamma)$ transitions is presented in Fig. 5. In difference

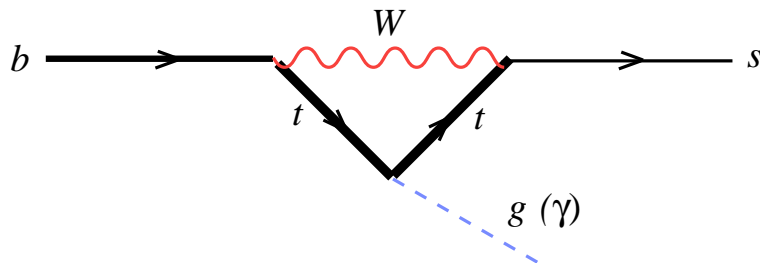


Figure 5: Magnetic penguins for $b \rightarrow sg(\gamma)$ transitions

with strange particle decays there is no suppression due to a small quark mass (m_b enters instead of m_s). From the diagram is seen that “magnetic” penguins are sensitive to heavy quarks in the loop and, for this reason, to possible deviations from the Standard Model prediction,

$$\text{Br}(b \rightarrow s\gamma) = (3.5 \pm 0.3) \cdot 10^{-4} \quad (26)$$

The inclusive rate of $B \rightarrow X_s\gamma$ decays was measured by CLEO and ALEPH collaborations,

$$\text{CLEO: } \text{Br}(B \rightarrow X_s\gamma) = (2.32 \pm 0.57 \pm 0.35) \cdot 10^{-4}$$

$$\text{ALEPH: } (3.11 \pm 0.80 \pm 0.72) \cdot 10^{-4}$$

No deviation was observed. This CLEO experiment was recognized at this meeting, Ed Thorndike received the 1999 Panofsky Prize.

4.3 Direct CP violation in K decays

At this meeting Robert Tschirhart reported a new measurement of direct CP violation in K decays by KTeV collaboration¹⁷

$$\frac{\epsilon'}{\epsilon} = (2.8 \pm 0.41) \cdot 10^{-3}, \quad (27)$$

which is in agreement with the CERN NA31 experiment but different from the previous measurement in the Fermilab E731 experiment.

The direct CP violation is a crucial test of the Standard Model. Does the result fit Standard Model? In his presentation Tschirhart defined an answer to this question as debatable. Indeed, in the recent work¹⁸ authors claim that the Standard Model leads to the predictions substantially (few times) lower than the experimental result (27), and only for extreme values of input parameters the theory can be consistent with the data.

My statement on the issue is that our theory of nonleptonic decays naturally predicts ϵ'/ϵ consistent with Eq. (27). Leaving detail for discussion elsewhere¹⁹, let me make few remarks.

Actually, relatively large values of ϵ'/ϵ within such approach was obtained long ago by Gilman and Wise²⁰, and by Voloshin²¹, but then the approach was unfairly abandoned. The argumentation was, see e.g. Ref.¹⁸, that the operator \mathcal{O}_5 enters with the small coefficient at $\mu = m_c$ or $\mu = 1\text{ GeV}$, so other operators are important. This criticism, however, is not relevant to the dominance of the operator \mathcal{O}_5 in the *low* normalization point where factorization takes place.

In one loop approximation direct CP violation shows up as an imaginary part of the coefficient c_5

$$\text{Im } c_5(m_c) = \frac{\text{Im}(V_{cs}^* V_{cd})}{V_{us}^* V_{ud}} \frac{\alpha_S(m_W)}{12\pi} \ln \frac{m_W^2}{m_c^2} \approx 0.12 \text{Im}(V_{cs}^* V_{cd})$$

due to the diagram of Fig. 1 with the c quark in the loop (t quark contribution is small). In difference with CP even part, i.e. $\text{Re } c_5$, which comes from virtual momenta between μ and m_c (thanks to the GIM cancellation), the CP odd part, $\text{Im } c_5$, comes from a larger range between m_c and m_W . The corrections due to logarithmic evolution can be simply accounted for, they increase the value of $\text{Im } c_5(m_c)$. Additionally $\text{Im } c_5(\mu)$ is increased by an evolution down to the normalization μ where $\alpha(\mu) \sim 1$ and factorization is applied.

Accounting for the phenomenological value (23) of CP even part we find ϵ'/ϵ in ballpark of the data. Besides confirmation of the quark mixing nature of CP violation in the Standard Model this serves, somewhat surprisingly, as one more confirmation of our mechanism for $\Delta I = 1/2$ enhancement. It is nice to get such a surprise on the eve of the Sakurai Prize.

5 Conclusions

Summarizing the development of twenty four years I cannot, unfortunately, say that theoretical understanding of $\Delta I = 1/2$ is very much advanced. Progress of last years in technically difficult calculations of higher loops corrections to OPE coefficients is not crucial when the effect comes from large numbers in matrix elements. As we discussed above the enhancement of $\Delta I = 1/2$ reflects the

existence of the large momentum scale in light hadrons. In case of glueballs the scale could be even larger, as it discussed in the next talk by Valya Zakharov²³. Moreover, the momentum scale in light hadrons related with this enhancement is so large that treatment of the c quark as heavy becomes questionable. In this sense the c quark is not heavy enough (let me remind, in passing, that a low upper limit for its mass was theoretically found²² in the Standard Model before the experimental discovery of the c quark).

Thus, we still have some distance to go so I finish by the sentence:

Penguins spread out but have not landed yet.

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References

1. J. Schwinger, Phys. Rev. Lett. **12**, 630 (1964).
2. K. Wilson, Phys. Rev. **179**, 1499 (1969).
3. A.Z. Patashinskii and V.L. Pokrovskii, *Fluctuation Theory of Phase Transitions*, Pergamon Press, 1979.
4. M.K. Gaillard and B. Lee, Phys. Rev. Lett. **33**, 108 (1974).
5. G. Altarelli and L. Miani, Phys. Lett. **B52**, 351 (1974).
6. S.L. Glashow, J. Iliopoulos, and L. Miani, Phys. Rev. **D2**, 1285 (1970).
7. H. Leutwyler, Phys. Lett. **B48**, 45 (1974); Nucl. Phys. **B76**, 413 (1974).
8. M. Gell-Mann, Elementary Particles, Oppenheimer Lectures, Institute for Advanced Study preprints 75-0978, 75-0979.
9. A.I. Vainshtein, V.I. Zakharov, and M.A. Shifman, Pisma Zh. Eksp. Teor. Fiz. **22**, 123 (1975) [JETP Lett. **22** 55 (1975)]; Nucl. Phys. **B120**, 316 (1977); Zh. Eksp. Teor. Fiz. **72**, 1275 (1977) [Sov. Phys. JETP **45** 670, (1977)].
10. J. Ellis, M.K. Gaillard, D.V. Nanopoulos, and S. Rudaz, Nucl. Phys. **B131**, 285 (1977), (E) **B132** 541 (1978)
11. M. Shifman, *ITEP Lectures in Particle Physics*, hep-ph/9510397.
12. F.J. Gilman and M.B. Wise, Phys. Rev. **D20**, 2392 (1979).

13. G. Buchalla, A.J. Buras, and M.E. Lautenbacher, *Rev. Mod. Phys.* **68**, 1125 (1996).
14. J.C. Pati and C.H. Woo, *Phys. Rev.* **D3**, 2920 (1972).
15. W.A. Bardeen, A.J. Buras, and J.-M. Gerard, *Phys. Lett.* **B192**, 138 (1987); *Nucl. Phys.* **B293**, 787 (1987).
16. K. Lingel, T. Skwarnicki, and J.G. Smith, *Penguin Decays of B Mesons*. hep-ex/9804015.
17. KTeV collaboration, *Observation of Direct CP Violation in $K_{S,L}$ to $\pi\pi$ Decays*, hep-ex/9905060.
18. S. Bosch *et al*, *Standard Model Confronting New Results for ep - $silon'/epsilon$* , hep-ph/9904408.
19. A. Vainshtein and M. Voloshin, in preparation.
20. F.J. Gilman and M.B. Wise, *Phys. Lett.* **83B**, 83 (1979).
21. M.B. Voloshin, *Introduction To Quantum Chromodynamics. Effects Of Strong Interactions In Nonleptonic Weak Decays*, Preprint ITEP-22-1981.
22. A.I. Vainshtein and I.B. Khriplovich, *Pisma Zh. Eksp. Teor. Fiz.* **18**, 141 (1973);
M.K. Gaillard and B.W. Lee, *Phys. Rev.* **D10**, 897 (1974)
23. V.I. Zakharov, *Gluon Condensate and Beyond*, The 1999 Sakurai Prize Lecture at the Centennial Meeting of the APS, hep-ph/9906xxx .